Novel mechanism of high-p_T production from an opaque quark-gluon plasma

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(Dated: February 8, 2008)

We find that in an opaque quark-gluon plasma, significant particle production at high transverse momenta can occur via acceleration of (initially) lower-momentum partons. This new mechanism, which is the opposite of parton energy loss (jet quenching), has important implications for observables, such as elliptic flow, in heavy-ion collisions at RHIC energies and above.

PACS numbers: 12.38.Mh, 25.75.-q, 25.75.Ld

Introduction. - There are several indications that an opaque, strongly-interacting quark-gluon plasma has been created [1, 2, 3] in relativistic heavy-ion collisions at the Relativistic Heavy Ion Collider (RHIC). One key observable is the large azimuthal momentum anisotropy of particle production [4, 5, 6] in noncentral collisions, characterized by the elliptic flow coefficient $v_2 \equiv \langle \cos 2\phi \rangle$, which requires more than an order of magnitude higher opacities than that of a weakly-coupled perturbative parton plasma [7]. Though such a dense system is still dissipative [8], i.e., it is *not* an ideal fluid, transport mean free paths $\lambda \sim 0.1$ fm are as short as their quantum mechanical lower bound [9] and, therefore, the system seems to be the most perfect fluid possible.

The large parton opacities have several consequences. For example, charm quarks and antiquarks are predicted to acquire a large elliptic flow [10], and space-momentum correlations of energetic particles are expected to show a strong surface emission pattern [11]. In this paper we show that particle production at high-p_T is also affected because *there is a considerable probability that a low-momentum parton gets accelerated to high transverse momenta via multiple scatterings.*

We emphasize that the effect considered here is different from the parton energy gain due to detailed balance in a *static* thermal system [12]. The new source of high-p_T particles here is the collective *dynamics* of the dense parton medium. These particles would be present in the final state, even if there were no energetic partons in the initial state.

Covariant parton transport theory is the incoherent particle limit of the underlying nonequilibrium quantum theory. Its main advantages are that i) it treats both the low-momentum bulk sector and high-momentum rare particles in a self-consistent framework; ii) it is not limited to local equilibrium; and iii) it describes the break-up of the system self-consistently.

We consider here an inelastic extension [10] of the transport equations in Refs. [7, 13, 14, 15], in which the on-shell parton phase space densities $\{f_i(x,\vec{p})\}$ evolve with elastic $2\to 2$ and *inelastic* $2\to 2$ rates as

$$p_1^\mu \partial_\mu f_{1,i} = S_i(x,\vec{p}_1) + \frac{1}{16\pi^2}\sum_{jk\ell}\iiint\limits_{2\,3\,4}\left(f_{3,k}f_{4,\ell}\frac{g_i g_j}{g_k g_\ell} - f_{1,i}f_{2,j}\right)\left|\bar{\mathcal{M}}^{ij\to k\ell}_{12\to 34}\right|^2 \delta^4(p_1+p_2-p_3-p_4)\ . \tag{1}$$

$|\bar{\mathcal{M}}|^2$ is the polarization averaged scattering matrix element squared, the integrals are shorthands for $\int\limits_a \equiv \int d^3p_a/(2E_a)$, g_i is the number of internal degrees of freedom for species i, while $f_{a,i} \equiv f_i(x,\vec{p}_a)$. The source functions $\{S_i(x,\vec{p})\}$ specify the initial conditions.

To study $Au+Au$ collisions at RHIC at $\sqrt{s_{NN}} = 200$ GeV with impact parameter $b = 8$ fm ($\approx 30\%$ central), we apply (1) to a system of massless gluons ($g = 16$), massless light (u,d) and strange quarks/antiquarks, and charm quarks/antiquarks ($g = 6$) with mass $M_c = 1.2$ GeV. The initial conditions and matrix elements were the same as in [10]. To model the strongly-interacting parton system at RHIC [7], parton densities were increased two-fold and parton cross sections were increased about three-fold to $dN(b{=}0)/d\eta = 2000$ at midrapidity and $\sigma_{gg\to gg} = 10$ mb, from their perturbative estimates $dN/d\eta \approx 1000$ and $\sigma_{gg\to gg} \approx 3$ mb. All elastic and inelastic leading-order $2\to 2$ QCD processes were taken into account and were assumed to be enhanced by the same factor 10/3. The initial momentum distributions were taken from leading-order perturbative QCD (with a K-factor of 2, GRV98LO PDFs, and $Q^2{=}p_T^2$, while $Q^2 = \hat{s}$ for charm), and the low-p_T divergence in the light-parton jet cross sections regulated via a smooth extrapolation below $p_\perp < 2$ GeV. The transverse density distribution was proportional to the binary collision dis-

tribution for two Woods-Saxon distributions, therefore $dN^{parton}(b{=}8\text{ fm})/dy \approx 500$. Perfect $\eta = y$ correlation was assumed.

The transport solutions were obtained via Molnar's Parton Cascade algorithm [16] (MPC), which employs the parton subdivision technique [17] to maintain Lorentz covariance and causality.

Results. - There are three general ways a particle can acquire its *final* transverse momentum during the transport evolution:

i) it has no interactions at all, in which case the particle comes from the *corona* of the collision region;

ii) it *loses* energy via interactions and moves to lower p_T, which is the jet quenching component; or

iii) it interacts and *gains* energy, i.e., gets pushed to higher p_T by the medium.

Figure 1 shows the relative fractions of these three components as a function of the final p_T from the transport calculation for partons at midrapidity ($|y_f| < 1$). At low final $p_T \sim 1-2$ GeV, partons predominantly arrive via the collective "push" from the medium. Contributions from quenched partons are smaller, and the corona is negligible. On the other hand, as p_T increases, the quenched component and the corona become more and more significant, while the accelerated part gradually decreases. However, even at $p_T \sim 7-8$ GeV, energy loss gives only one half of the yield, while almost *one-third* of the partons are accelerated (initially) lower-momentum partons.

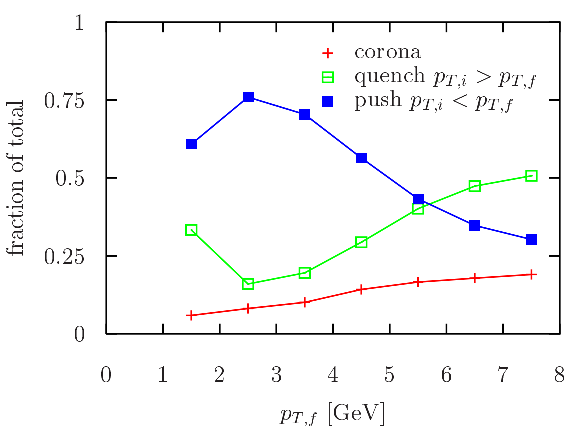

FIG. 1: Fraction of the final high-p_T yield coming from partons without interactions (corona), from energy loss (quench), and from accelerated partons (push) as a function of p_T in $Au + Au$ at $\sqrt{s_{NN}} = 200$ GeV with $b = 8$ fm, for ≈ 6 times higher opacities than perturbative estimates. Computed via MPC [7, 16].

Figure 2 gives a more detailed picture, where the normalized initial p_T distributions are shown for partons in various final p_T intervals (and final rapidity $|y_f| < 1$), with contributions from the corona omitted. For all final p_T intervals, the accelerated soft partons predominantly come from $p_{T,i} \sim 1$ GeV, which is precisely where the initial parton p_T distribution (dashed line) has a strong peak. At the opacities considered here, partons have $\langle N_{coll}\rangle \approx 14$ scatterings on average. Though the likelihood of gaining significant p_T, even in so many scatterings, is small, the initial high yield of low-momentum partons multiplied with the low probability can still compete with the much lower parton yields at high p_T. As the final p_T increases, the relative importance of low-momentum partons decreases, in accordance with the decrease in the fractional "push" yield in Fig. 1.

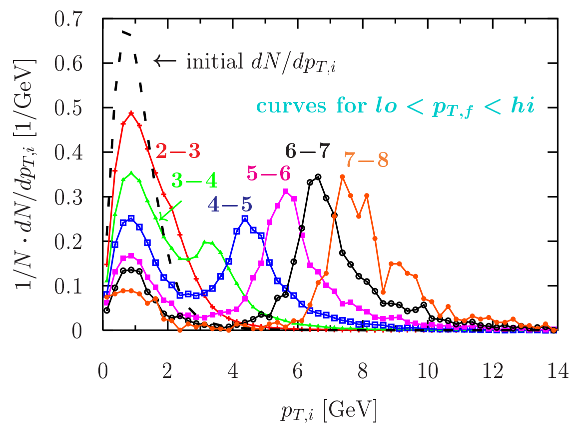

FIG. 2: Normalized initial p_T distributions for partons in various final p_T intervals (and final $|y_f| < 1$), omitting contributions from the corona, from the same calculation as Fig. 1. The initial parton p_T distribution is also shown (dashed line).

Two-component models based on a thermal (or hydrodynamic) collective part and a perturbative QCD jet part underestimate the collective component at high p_T because the approximately exponential p_T distribution imposed (with typical few-hundred-MeV slope parameters) cannot compete with the power-law jet spectrum. However, such a characterization of the collective part is based on macroscopic approaches that are formulated in terms of the lowest few moments (e.g., particle density, energy density) of the local momentum distribution, and therefore cannot be trusted for the high-momentum tails. Our results in Figs. 1 and 2 imply that (initially) low-momentum partons that suffer *finite* number of scatterings populate high transverse momenta with a non-exponential, power-law-like distribution. In fact, already a single scattering generates power-law tails because for typical Debye-screened t-channel gluon exchange $d\sigma/dt \sim \alpha_s^2/(t-\mu_D^2)^2 \sim \alpha_s^2/(\vec{q}^{\,2}+\mu_D^2)^2$, where $\vec{q}$ is the transferred momentum.

The above findings have significant implications for elliptic flow. Figure 3 shows elliptic flow as a function of p_T for the three components (corona, quench, push) and for their combined yield. For partons that have lost energy (quench), elliptic flow drops rather sharply at high p_T, very similarly to calculations based on inelastic pertur-

bative QCD energy loss in the Eikonal limit [18]. In the transport calculation, however, the sharp drop is largely compensated by the much larger v_2 of the accelerated partons (push). The significant "push" component also means that the "geometric" elliptic flow bounds argued in [19], which assumed extreme quenching, do not apply. Though the geometric corona alone gives a positive, p_T-independent elliptic flow (reflecting the initial spatial anisotropy of the edges of the collision zone), it is much smaller than the contribution coming from the accelerated partons. Therefore, the collective dynamics of the opaque quark-gluon plasma can provide a natural explanation for why elliptic flow at RHIC exceeds the "geometric" bounds and why it decreases relatively slowly at high transverse momenta [20].

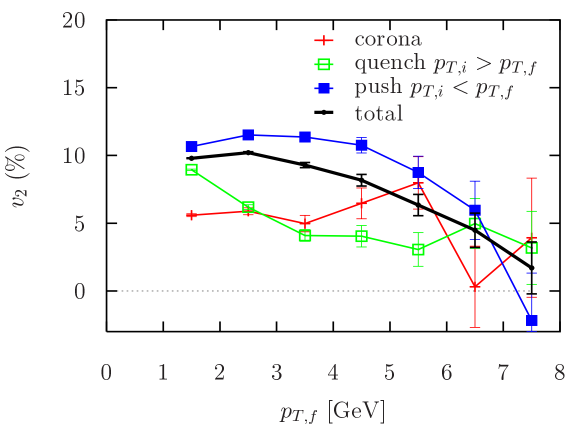

FIG. 3: Elliptic flow as a function of p_T for partons without interactions (corona), from energy loss (quench), from accelerated partons (push), and from all three components combined (total), from the same calculation as Fig. 1.

Conclusions. - We show from covariant parton transport theory that in an expanding opaque quark-gluon plasma there is a considerable probability that low-momentum partons get accelerated to high transverse momenta in multiple scatterings. This new mechanism of high-p_T production, which is the opposite of jet quenching, has significant implications for observables at RHIC energies and beyond. For example, it gives a natural explanation to why elliptic flow at RHIC exceeds so called "geometric" bounds [19] and decreases relatively slowly at high transverse momenta.

This study considered six times higher opacities than that of a weakly-coupled perturbative parton plasma, which is 2-3 times smaller than the opacities estimated from RHIC data [7] and likely even a smaller fraction of the opacities reachable at the future Large Hadron Collider (LHC). At larger opacities, accelerated partons would have even larger importance and would influence yields out to higher $p_T > 7 - 8$ GeV. In addition, this work investigated only $2 \to 2$ elastic and inelastic transport. A more complete study will also have to consider in the future the effect of radiative $2 \leftrightarrow 3$ processes.

Acknowledgments. - Discussions with M. Gyulassy are acknowledged. This work was supported by DOE grant DE-FG02-01ER41190.

[1] M. Gyulassy and L. McLerran, Nucl. Phys. A **750**, 30 (2005).
[2] E. V. Shuryak, hep-ph/0405066.
[3] H. Stoecker, Nucl. Phys. A **750**, 121 (2005).
[4] C. Adler *et al.* [STAR Collaboration], Phys. Rev. Lett. **90**, 032301 (2003); K. Filimonov [STAR Collaboration], Nucl. Phys. A**715**, 737 (2003).
[5] K. Adcox *et al.* [PHENIX Collaboration], Phys. Rev. Lett. **89**, 212301 (2002); N. N. Ajitanand [PHENIX Collaboration], Nucl. Phys. A **715**, 765 (2003).
[6] B. B. Back *et al.* [Phobos Collaboration], nucl-ex/0407012.
[7] D. Molnar and M. Gyulassy, Nucl. Phys. **A697**, 495 (2002), **A703**, 893(E) (2002); *ibid.* **A698**, 379 (2002).
[8] D. Molnar and P. Huovinen, Phys. Rev. Lett. **94**, 012302 (2005).
[9] P. Danielewicz and M. Gyulassy, Phys. Rev. D **31**, 53 (1985).
[10] D. Molnar, nucl-th/0410041.
[11] D. Molnar, nucl-th/0408044.
[12] E. Wang and X. N. Wang, Nucl. Phys. A **715**, 787 (2003); J. A. Osborne, E. Wang and X. N. Wang, Phys. Rev. D **67**, 094022 (2003).
[13] B. Zhang, M. Gyulassy and C. M. Ko, Phys. Lett. B **455**, 45 (1999).
[14] B. Zhang, Comput. Phys. Commun. **109**, 193 (1998).
[15] D. Molnar and M. Gyulassy, Phys. Rev. C **62**, 054907 (2000).
[16] D. Molnár, MPC 1.8.0. This parton cascade code is available at http://nt3.phys.columbia.edu/people/molnard.
[17] Pang Y 1996 *Preprint* RHIC 96 Summer Study, CU-TP-815 (unpublished)
[18] X.-N. Wang, *Phys. Rev.* C **63**, 054902 (2001); M. Gyulassy M, I. Vitev and X.-N. Wang, *Phys. Rev. Lett.* **86**, 2537 (2001); M. Gyulassy *et al. Phys. Lett.* B **526**, 301 (2002).
[19] E. V. Shuryak, Phys. Rev. C **66**, 027902 (2002).
[20] J. Adams *et al.* [STAR Collaboration], Phys. Rev. Lett. **93**, 252301 (2004).